

Identification of a Scalar Glueball

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We have performed a coupled channel study of the meson-meson S-waves involving isospins (I) 0, 1/2 and 3/2 up to 2 GeV. For the first time the channels $\pi\pi$, $K\bar{K}$, $\eta\eta$, $\sigma\sigma$, $\eta\eta'$, $\eta'\eta'$, $\rho\rho$, $\omega\omega$, $\omega\phi$, $\phi\phi$, $a_1\pi$ and $\pi^*\pi$ are considered. All the resonances with masses below 2 GeV for $I = 0$ and 1/2 are generated by the approach. We identify the $f_0(1710)$ and a pole at 1.6 GeV, which is an important contribution to the $f_0(1500)$, as glueballs. This is based on an accurate agreement of our results with predictions of lattice QCD and the chiral suppression of the coupling of a scalar glueball to $\bar{q}q$. Another nearby pole, mainly corresponding to the $f_0(1370)$, is a pure octet state not mixed with the glueball.

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1. QCD, the present theory of strong interactions, is a non-abelian Yang Mills theory so that gluons carry colour charge and interact between them. It is generally believed that QCD predicts the existence of mesons without valence quarks, the so called glueballs. Its confirmation in the spectrum of strong interactions is then at the heart of the theory. In quenched lattice QCD the lightest glueball has the quantum numbers of the vacuum, $J^{PC} = 0^{++}$, with a mass of (1.66 ± 0.05) GeV [1]. Experimentally the closest 0^{++} scalar resonances to this energy range are the $f_0(1500)$ and $f_0(1710)$ [2]. Some references favour the $f_0(1500)$ as the lightest scalar glueball [3], while others do so for the $f_0(1710)$ [4,5].

We analyze the $I = 0$ meson-meson S-wave in terms of 13 coupled channels, $\pi\pi(1)$, $K\bar{K}(2)$, $\eta\eta(3)$, $\sigma\sigma(4)$, $\eta\eta'(5)$, $\eta'\eta'(6)$, $\rho\rho(7)$, $\omega\omega(8)$, $K^*\bar{K}^*(9)$, $\omega\phi(10)$, $\phi\phi(11)$, $a_1(1260)\pi(12)$ and $\pi^*(1300)\pi(13)$. The number labelling each state is given between brackets to the right. The multipion states, which play an increasing role for energies above ~ 1.2 GeV, are mimicked through the $\sigma\sigma$, $\rho\rho$ and $\omega\omega$ channels. It is worth stressing that our approach is the first one with such a large number of channels and that a similar scheme could also be applied to other controversial meson-meson partial waves. In addition, we study simultaneously the S-wave of $K^-\pi^+ \rightarrow K^-\pi^+$ (involving $I = 1/2$ and $3/2$) with the coupled channels $K\pi$, $K\eta$ and $K\eta'$.

2. Let $T_{i,j}^{(I)}$ be the $i \leftrightarrow j$ S-wave amplitude with isospin I and $i, j = 1 \dots n$, with n the number of channels. We use the master formula $T^{(I)} = [I + N \cdot g]^{-1} \cdot N$, where N is the symmetric matrix of interaction kernels and g is a diagonal matrix of elements $g_i(s)$. The function $g_i(s)$ is calculated from kinematics in terms of a once subtracted dispersion relation and a subtraction constant a_i [6]. Since SU(3) breaking is milder in the vector sector we take $a_7 = a_8 = a_9 = a_{10} = a_{11}$. The rest of subtraction constants are fitted to data. The matrix elements $N_{i,j}$ consist of the sum of two tree level contributions. The first is a contact interaction calcu-

lated from the lowest order Chiral Perturbation Theory Lagrangian, \mathcal{L}_2 . The second is due to the exchange of bare resonances in the s -channel with the couplings calculated from the lowest order chiral Lagrangian including an octet and singlet of 0^{++} resonances, \mathcal{L}_S [7]. Explicit expressions of $N_{i,j}$ can be found in ref. [6] for the simplified case of three channels without including the η_1 field. We extend these Lagrangians from SU(3) to U(3) as the η_1 field is needed to deal with the η and η' mesons, similarly as in ref. [8]. The matrix $\Phi = \sum_{i=1}^8 \phi_i \lambda_i / \sqrt{2} + \eta_1 / \sqrt{3}$ incorporates in a standard way the nonet of the lightest pseudoscalars. We also employ the matrix $U = \exp(i\sqrt{2}\Phi/f)$ and the covariant derivative $D_\mu U = \partial_\mu U - ir_\mu U + iU\ell_\mu$, with f the pion decay constant in the chiral limit fixed to $f_\pi = 92.4$ MeV. The classical left and right external fields, r_μ and ℓ_μ , respectively, are necessary to gauge the global chiral symmetry to a local one [7]. The field $v_\mu = (r_\mu + \ell_\mu)/2$ plays a special role in our approach since it is identified with λW_μ , where W_μ is the nonet of the lightest 1^{--} vector resonances and λ is a constant, with $\lambda = 4.3$ from the width $\rho \rightarrow \pi\pi$. The couplings of the vector-vector states to the pseudoscalar-pseudoscalar and $\sigma\sigma$ ones are then determined by minimal coupling [9]. Our fits require a singlet and two octets of bare resonances. The two octets were already considered in ref. [8] in the study of $K^-\pi^+ \rightarrow K^-\pi^+$. We fix the parameters of the first octet, mass and coupling constants, to those in ref. [8], $M_8^{(1)} = 1.29$ GeV, $c_d^{(1)} = c_m^{(1)} = 26$ MeV. The bare mass of the second octet is fixed from the same reference, $M_8^{(2)} = 1.90$ GeV. We are then left with three parameters for the singlet, M_1 , $\tilde{c}_d^{(1)}$, $\tilde{c}_m^{(1)}$, and two for the second octet, $c_d^{(2)}$ and $c_m^{(2)}$. It results from our fits that $M_1 \lesssim 0.9$ GeV.

Concerning the $\sigma\sigma$ channel we follow a novel method to calculate its transition amplitudes, $N_{i,4}$, without including any new free parameter. This can be done because the σ corresponds to a pole due to the interactions between two pions in the $I = 0$ S-wave, $(\pi\pi)_0$

[10]. For the interaction kernel $N_{i,4}$ one starts by calculating from the Lagrangians \mathcal{L}_2 and \mathcal{L}_S the tree level amplitude $T_{i,4}^{2+S}$ for $i \rightarrow (\pi\pi)_0(\pi\pi)_0$. To take into account the pion final state interactions, $T_{i,4}^{2+S}$ is multiplied by the factor $\prod_{k=1}^m 1/D(s_k)$, with m the number of σ 's in the scattering process (2 or 4) and s_k the total centre of mass (CM) energy squared of the k_{th} pair. We use here that the rescattering of two $I = 0$ S-wave pions from a production kernel is given by the factor $1/D = 1/(1 + V_2 g_1)$, with $V_2 = (s - m_\pi^2/2)/f^2$ [10]. To isolate $N_{i,4}$ one takes the limit (for definiteness $i \neq 4$)

$$\lim_{s_1, s_2 \rightarrow s_\sigma} \frac{T_{i,4}^{2+S}}{D_{II}(s_1)D_{II}(s_2)} = \frac{N_{i,4} g_{\sigma\pi\pi}^2}{(s_1 - s_\sigma)(s_2 - s_\sigma)}. \quad (1)$$

Where the subscript II indicates that the corresponding function is calculated on the second Riemann sheet (with the sign reversed in the definition of the pion three-momentum), s_σ is the σ pole position and $g_{\sigma\pi\pi}$ is its coupling to $\pi\pi$. Performing the Laurent expansion around s_σ of $1/D_{II}(s) = \alpha_0/(s - s_\sigma) + \dots$ the evaluation of $N_{i,4}$ from eq.(1) requires the ratio $(\alpha_0/g_{\sigma\pi\pi})^2$. Since $g_{1,II}(s_\sigma) = -f^2/(s_\sigma - m_\pi^2/2)$ at s_σ , where $1 + V_2 g_{1,II} = 0$, and taking $T_{II} \simeq V_2/(1 + V_2 g_{1,II})$, appropriate for these energies [10], then $(\alpha_0/g_{\sigma\pi\pi})^2 = f^2/(1 - \frac{dg_{1,II}}{ds}|_{s_\sigma} \frac{(s_\sigma - m_\pi^2/2)^2}{f^2}) \simeq f^2$. In this way, $N_{i,4} = T_{i,4}^{2+S} f^2$, $i \neq 4$, and $N_{4,4} = T_{4,4}^{2+S} f^4$. Using $N_{i,4}$ evaluated with $s_k = s_\sigma$ violates unitarity because s_σ is complex and $N_{i,4}$ must be real. Instead, we interpret the width of the σ resonance as a Lorentzian mass distribution around its nominal mass value ~ 450 MeV with a width ~ 500 MeV. In this way the σ masses ($\sqrt{s_k}$) used to calculate the functions $N_{i,4}$ and g_4 are folded with the previous mass distribution. Similarly, for the $\rho\rho$ state g_7 is also convoluted with a ρ mass distribution.

3. We fit our 12 free parameters to 370 data points from threshold up to 2 GeV. The data comprise the $I = 0$ S-wave $\pi\pi$ phase shifts δ_0^0 , the elasticity $\eta_0^0 = |S_{1,1}|$, the $I = 0$ S-wave $\pi\pi \rightarrow K\bar{K}$ phase shifts $\delta_{1,2}$ and modulus $|S_{1,2}|$, the S-wave contribution to the $\pi\pi \rightarrow \eta\eta, \eta\eta'$ event distributions and the phase (ϕ) and modulus (A) of the $K^-\pi^+ \rightarrow K^-\pi^+$ amplitude from the LASS data. The S-matrix element $S_{i,j}$ is given by $S_{i,j} = \delta_{ij} + 2i\sqrt{\rho_i} T_{i,j}^{(I)} \sqrt{\rho_j}$, where $\rho_i = q_i/8\pi\sqrt{s}$ and q_i is the CM three-momentum for channel i . In order, these data are shown on the first eight panels of Fig.1 from top to bottom and left to right. For $\sqrt{s} \leq m_K$ in the $\delta_0^0(s)$ panel we have the inset showing in detail the precise data from K_{e4} decays. The reproduction of the data is fair, as shown in the figure. The dashed lines on the first eight panels include the $a_1\pi$ and $\pi^*\pi$ states, while the solid ones do not. The similarity between both curves indicates that these channels give

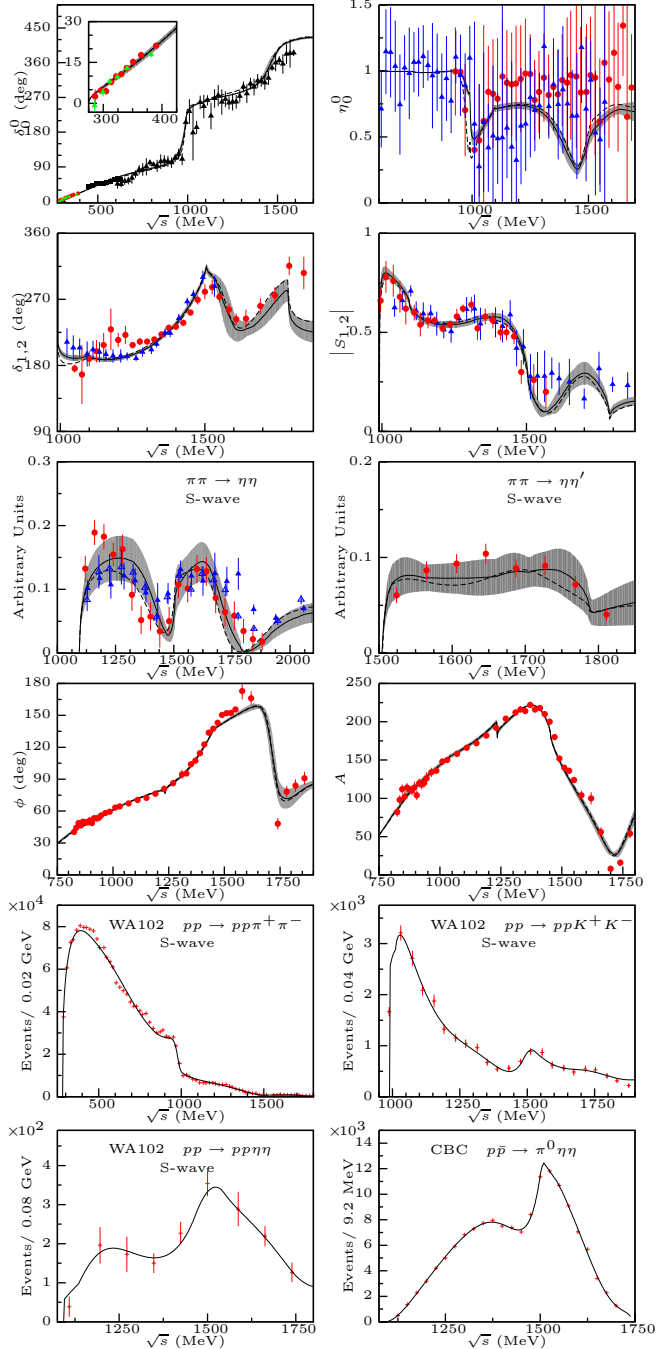


FIG. 1. Fit to experimental data. More details are given in the text.

small contributions. The width of the band represents our systematic uncertainties at the level of two standard deviations, $n_\sigma = \Delta\chi^2/(2\chi^2)^{1/2}$ [11]. Compared with other works [12–14] we determine the interaction kernels from standard chiral Lagrangians, avoid ad-hoc parameterizations, include many more channels and fewer free parameters are used. For $I = 1/2$ the κ pole is

located at $(708 \pm 6 - i 313 \pm 10)$ MeV, the $K_0^*(1430)$ at $(1435 \pm 6 - i 142 \pm 8)$ MeV and the $K_0^*(1950)$ at $(1750 \pm 20 - i 150 \pm 20)$ MeV, similarly to ref. [8]. For $I = 0$ one has the $f_0(600)$ or σ at $(456 \pm 6 - i 241 \pm 7)$ MeV and the $f_0(980)$ at $(983 \pm 4 - i 25 \pm 3)$ MeV. There are poles at $(1690 \pm 20 - i 110 \pm 20)$ MeV, corresponding to the $f_0(1710)$, and at $(1810 \pm 15 - i 190 \pm 20)$ MeV, with mass and width in agreement with those reported for the $f_0(1790)$ by BESII. In the PDG [2] the width for the $f_0(1710)$ is 137 ± 8 MeV, much smaller than 220 ± 40 MeV from the given pole position. However, we have checked that on the real axis the value of the width corresponding to the half-maximum for the partial waves with prominent $f_0(1710)$ peaks is just 160 MeV [15]. This reduction is due to the opening of several channels along the resonance region and the agreement with the PDG is restored. The other poles at $(1466 \pm 15 - i 158 \pm 12)$ MeV and $(1602 \pm 15 - i 44 \pm 15)$ MeV, connected with the $f_0(1370)$ and $f_0(1500)$, are referred in the following as f_0^L and f_0^R , respectively. Despite that we have included only three bare resonances in $I = 0$ we have generated six. The poles are located on the unphysical Riemann sheets that connect continuously with the physical one for some interval along the real s -axis. Note that the pole f_0^R does not influence the physical axis beyond the $\eta\eta'$ threshold at 1505 MeV, since above this energy a different Riemann sheet is the one that matches with the physical s -axis. This effect typically gives rise to a pronounced signal at the $\eta\eta'$ threshold and this is the reason for the $f_0(1500)$ mass, (1505 ± 6) MeV [2]. If a physical amplitude is dominated by the f_0^R pole, then its peak at 1505 MeV has an effective width larger than the one from the pole position, 88 MeV. This is so because given a Breit-Wigner located at the position of the f_0^R pole the energy interval below 1.5 GeV at which half the value of the modulus squared at 1.5 GeV is reached is $\delta = 1.2\Gamma = 105$ MeV, the width of the $f_0(1500)$ [2]. The $f_0(1370)$ is mainly given by the f_0^L pole, though its precise shape is sensitive to f_0^R for those channels that couple strongly with the latter. In Fig.1 we also show in the last two rows data from pp inelastic scattering at 450 GeV/c and $p\bar{p}$ annihilation by the WA102 and Crystal Barrel (CBC) Collaborations, respectively. We have fitted the WA102 data using a coherent sum of Breit-Wigner functions and a non-resonant term, similarly as done by the WA102 Collaboration [16]:

$$i)\sqrt{s} < m_\eta + m_{\eta'}, A = \{\sigma, f_0(980), f_0^L, f_0^R\},$$

$$A(\sqrt{s})_i = NR(\sqrt{s})_i + \sum_{j \in A} \frac{a_j e^{i\theta_j} g_{j;i}}{M_j^2 - s - i M_j \Gamma_j},$$

$$ii)\sqrt{s} > m_\eta + m_{\eta'}, B = \{\sigma, f_0(980), f_0(1710), f_0(1790)\},$$

$$A(\sqrt{s})_i = NR(\sqrt{s})_i + r_i + \sum_{j \in B} \frac{a_j e^{i\theta_j} g_{j;i}}{M_j^2 - s - i M_j \Gamma_j},$$

$$NR(\sqrt{s})_i = \alpha(\sqrt{s} - m_k - m_\ell)^\beta e^{-\gamma\sqrt{s} - \delta s}, \quad (2)$$

where a_j and θ_j are the amplitude and the phase of the production vertex of the j th resonance, M_j , Γ_j and $g_{j;i}$ are, respectively, the mass, width and the coupling to channel i of the same resonance. The latter is determined from the residue of the partial waves at the pole position. In addition, $m_k + m_\ell$ is the threshold for the channel i and $\alpha, \beta, \gamma, \delta$ are real parameters. The form of the non-resonant term is taken from the WA102 Collaboration [16]. The constant r_i is fixed so as the amplitude $A(\sqrt{s})_i$ is continuous at $w_{\eta\eta'} \equiv m_\eta + m_{\eta'}$. As explained above, once the $\eta\eta'$ threshold is crossed over one has to consider other Riemann sheets which do not have the f_0^L and f_0^R poles but the $f_0(1710)$ and $f_0(1790)$ ones. Above $w_{\eta\eta'}$ the σ and $f_0(980)$ give tiny contributions. Γ_j in eq.(2) is the largest between its value from the pole position and the one calculated by summing the partial decay widths $\Gamma_{j;i} = \theta(\sqrt{s} - m_k - m_\ell) \lambda_i |g_{j;i}|^2 q_i / (8\pi M_j^2)$, with $\lambda_i = 1/2$ for identical particles. Eq.(2) incorporates important new facts compared to the analyses of the WA102 Collaboration. First, the pole positions for the different resonances are those already determined from our study of the scattering data on the first 8 panels of Fig.1. Let us stress that these observables only involve two particles in the final state and their analysis is theoretically cleaner. Second, the couplings $g_{j;i}$ are similarly fixed. Third, the a_j and θ_j parameters are the same for all the WA102 reactions, that are fitted simultaneously. For the Crystal Barrel data on $p\bar{p}$ annihilation we also use eq.(2) but without $NR(\sqrt{s})$. A good reproduction of the data results. In $p\bar{p} \rightarrow \pi^0 \eta\eta$ one observes a broad bump for the $f_0(1370)$ and a prominent peak for the $f_0(1500)$, that also gives strong signals in the WA102 data. Other peaks are observed for the σ , $f_0(980)$ and $f_0(1710)$. The latter is important for the the shoulder in $pp \rightarrow pp\eta\eta$ above 1.5 GeV.

GeV	$f_0(1370)$	f_0^R	$f_0(1710)$
$ g_{\pi^+\pi^-} $	3.59 ± 0.16	1.30 ± 0.22	1.21 ± 0.16
$ g_{K^0\bar{K}^0} $	2.23 ± 0.18	2.06 ± 0.17	2.0 ± 0.3
$ g_{\eta\eta} $	1.7 ± 0.3	3.78 ± 0.26	3.3 ± 0.8
$ g_{\eta\eta'} $	4.0 ± 0.3	4.99 ± 0.24	5.1 ± 0.8
$ g_{\eta'\eta'} $	3.7 ± 0.4	8.3 ± 0.6	11.7 ± 1.6

TABLE I. Couplings of the $f_0(1370)$, f_0^R and $f_0(1710)$.

4. In table I we give the couplings of the f_0^L (iden-

tified as the $f_0(1370)$), f_0^R and $f_0(1710)$ poles to the two pseudoscalar channels. We observe that the couplings of the f_0^R and $f_0(1710)$ are quite similar. This is so because the two poles coalesce in the same one when moving continuously from the sheet of one of them to the one of the other. They correspond to the same underlying resonance, but split in two due to the interaction in coupled channels. From the couplings of the $f_0(1710)$ one can calculate the branching ratios $\Gamma(K\bar{K})/\Gamma_{total} = 0.36 \pm 0.12(0.38^{+0.09}_{-0.19})$, $\Gamma(\eta\eta)/\Gamma_{total} = 0.22 \pm 0.12(0.18^{+0.03}_{-0.13})$, and $\Gamma(\pi\pi)/\Gamma(K\bar{K}) = 0.32 \pm 0.14(< 0.11)$, where the values of the PDG are given between brackets. The values are compatible within one sigma. We also obtain that the $f_0(1790)$ has a small $K\bar{K}$ coupling, and this is a major difference with respect to the $f_0(1710)$ as stressed by BESII. The couplings of the f_0^L - $f_0(1370)$ in table I correspond to the pure $I = 0$ octet member $(\bar{u}u + \bar{d}d - 2\bar{s}s)/\sqrt{6}$ because they are very close to the tree level ones $|g_{\pi^+\pi^-}| = 3.9$, $|g_{K^0\bar{K}^0}| = 2.3$, $|g_{\eta\eta}| = 1.4$, $|g_{\eta'\eta'}| = 3.7$, $|g_{\eta'\eta'}| = 3.8$ GeV calculated from the Lagrangian \mathcal{L}_S [7], with $c_d^{(1)}$, $c_m^{(1)}$ and $M_8^{(1)}$ given above. We have also checked that this is the case for the $K_0^*(1430)$ resonance which is the $I = 1/2$ member of the same octet. It follows then that the first octet is a pure one without mixing with the nearby f_0^R and $f_0(1710)$. The f_0^L - $f_0(1370)$ couplings imply a large width to $\pi\pi$ with $\Gamma(f_0(1370) \rightarrow 4\pi)/\Gamma(f_0(1370) \rightarrow \pi\pi) = 0.30 \pm 0.12$, in good agreement with the interval 0.10-0.25 given in the recent ref. [17]. Let us see that the pattern of sizes of the couplings of the f_0^R and $f_0(1710)$ corresponds to the chiral suppression of the coupling of a scalar glueball, G_0 , to $\bar{q}q$ [5]. According to ref. [5] this coupling is proportional to the quark mass, which then implies a strong suppression in the production of $\bar{u}u$ and $\bar{d}d$ relative to $\bar{s}s$ from G_0 . With a pseudoscalar mixing angle $\sin\beta = -1/3$ one has that $\eta = -\eta_s/\sqrt{3} + \eta_u\sqrt{2/3}$ and $\eta' = \eta_s\sqrt{2/3} + \eta_u/\sqrt{3}$ with $\eta_s = \bar{s}s$ and $\eta_u = (\bar{u}u + \bar{d}d)/\sqrt{2}$. Denoting by g_{ss} the production of $\eta_s\eta_s$, g_{sn} that of $\eta_s\eta_u$ and g_{nn} for $\eta_u\eta_u$,

$$\begin{aligned} g_{\eta'\eta'} &= 2g_{ss}/3 + g_{nn}/3 + 2\sqrt{2}g_{ns}/3, \\ g_{\eta\eta'} &= -\sqrt{2}g_{ss}/3 + \sqrt{2}g_{nn}/3 + g_{ns}/3, \\ g_{\eta\eta} &= g_{ss}/3 + 2g_{nn}/3 - 2\sqrt{2}g_{ns}/3. \end{aligned} \quad (3)$$

If the chiral suppression of ref. [5] operates then $|g_{ss}| \gg |g_{nn}|$. This together with the OZI rule suppress the coupling g_{ns} . Taking e.g. the couplings of f_0^R one obtains $g_{ss} = 11.5 \pm 0.5$, $g_{ns} = -0.2$ and $g_{nn} = -1.4$ GeV, and the strong suppression is clear. We now consider the $K\bar{K}$ coupling. A K^0 in terms of valence quarks corresponds to $\sum_{i=1}^3 \bar{s}_i u^i/\sqrt{3}$, summing over the colour indices, and analogously for the \bar{K}^0 . The production of a

colour singlet $\bar{s}s$ from the $K^0\bar{K}^0$ requires then the combination $\bar{s}_i s^j = \delta_i^j \bar{s}s/3 + (\bar{s}_i s^j - \delta_i^j \bar{s}s/3)$, and similarly for $\bar{u}_j u^i$. As the production occurs from the colour singlet $\bar{s}s$ source, only the configuration $\bar{s}s\bar{u}u$ contributes, picking up a suppression factor of 1/3. In addition, the coupling g_{ss} has an extra factor 2 compared to that of a $\bar{s}s\bar{u}u$, because the former contains two $\bar{s}s$. One then expects that the coupling to $K^0\bar{K}^0$ has the absolute value $g_{ss}/6$. For the f_0^R and $f_0(1710)$ it results $|g_{K^0\bar{K}^0}| \simeq 2$ GeV, in good agreement with table I. Another resonance with a known enhanced coupling to $\bar{s}s$ is the $f_0(980)$. However, the sizes of its couplings to $\eta\eta$, $\eta\eta'$ and $\eta'\eta'$ follow the opposite order to the $f_0(1790)$ and f_0^R cases and all of them are much smaller than the coupling to $K\bar{K}$. Note that quenched lattice QCD [4] establishes that the couplings of the lightest scalar glueball to pseudoscalar pairs in the SU(3) limit scales as the quark mass, in support of the chiral suppression mechanism of ref. [5], that we also observe as discussed above. This mechanism also implies that the glueball should remain unmixed. This accurately fits with our previous result that both the f_0^R and $f_0(1710)$ do not mix with the nearby f_0^L . In addition, the masses of the f_0^R and $f_0(1710)$ poles are in excellent agreement with the quenched latticed QCD prediction for the mass of the lightest glueball, (1.66 ± 0.05) GeV.

5. In summary, we have presented a coupled channel study of the $I = 0, 1/2$ meson-meson S-waves from $\pi\pi$ threshold up to 2 GeV with 13 coupled channels. All the $I = 0$ and $1/2$ 0^{++} resonances with masses below 2 GeV have been generated. The $f_0(1710)$ and a pole at 1.6 GeV, which is an important contribution to the $f_0(1500)$, are identified as glueballs. Another pole at $(1.466 - i0.157)$ GeV, mainly corresponding to the $f_0(1370)$, is shown to be a pure octet member.

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